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Realistic Type IIB Supersymmetric Minkowski Flux Vacua

Ching-Ming Chen,[1] Tianjun Li,[1,2] Yan Liu,[2] and Dimitri V. Nanopoulos[1,3,4]

[1] *George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, TX 77843, USA*

[2] *Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100080, China*

[3] *Astroparticle Physics Group, Houston Advanced Research Center (HARC), Mitchell Campus, Woodlands, TX 77381, USA*

[4] *Academy of Athens, Division of Natural Sciences, 28 Panepistimiou Avenue, Athens 10679, Greece*

We show that there exist supersymmetric Minkowski vacua on Type IIB toroidal orientifold with general flux compactifications where the RR tadpole cancellation conditions can be relaxed elegantly. Then we present a realistic Pati-Salam like model. At the string scale, the gauge symmetry can be broken down to the Standard Model (SM) gauge symmetry, the gauge coupling unification can be achieved naturally, and all the extra chiral exotic particles can be decoupled so that we have the supersymmetric SMs with/without SM singlet(s) below the string scale. The observed SM fermion masses and mixings can also be obtained. In addition, the unified gauge coupling, the dilaton, the complex structure moduli, the real parts of the Kähler moduli and the sum of the imaginary parts of the Kähler moduli can be determined as functions of the four-dimensional dilaton and fluxes, and can be estimated as well.

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Introduction – One of the great challenging and essential problems in string phenomenology is the construction of the realistic string vacua, which can give us the low energy supersymmetric Standard Models (SMs) without exotic particles, and can stabilize the moduli fields. With renormalization group equation running, we can connect such constructions to the low energy realistic particle physics which will be tested at the upcoming Large Hadron Collider (LHC). During the last a few years, the intersecting D-brane models on Type II orientifolds [1], where the chiral fermions arise from the intersections of D-branes in the internal space [2] and the T-dual description in terms of magnetized D-branes [3], have been particularly interesting [4].

On Type IIA orientifolds with intersecting D6-branes, many non-supersymmetric three-family Standard-like models and Grand Unified Theories (GUTs) were constructed in the beginning [5]. However, there generically existed uncancelled Neveu-Schwarz-Neveu-Schwarz (NSNS) tadpoles and the gauge hierarchy problem. To solve these problems, semi-realistic supersymmetric Standard-like and GUT models have been constructed in Type IIA theory on the $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [6, 7] and other backgrounds [8]. Interestingly, only the Pati-Salam like models can give all the Yukawa couplings. Without the flux background, Pati-Salam like models have been constructed systematically in Type IIA theory on the $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [7]. Although we may explain the SM fermion masses and mixings in one model [9], the moduli fields have not been stabilized, and it is very difficult to decouple the chiral exotic particles. To stabilize the moduli via supergravity fluxes, the flux models on Type II orientifolds have also been constructed [10, 11]. Especially, some flux models [11] can explain the SM fermion masses and mixings. However, those models are in the AdS vacua and have quite a few chiral exotic particles that are difficult to be decoupled.

In this paper, we consider the Type IIB toroidal orientifold with the Ramond-Ramond (RR), NSNS, non-geometric and S-dual flux compactifications [12]. We find that the RR tadpole cancellation conditions can be relaxed elegantly in the supersymmetric Minkowski vacua, and then we may construct the realistic Pati-Salam like models [13]. In this paper, we present a concrete simple model which is very interesting from the phenomenological point of view and might describe Nature. We emphasize that we do not fix the four-dimensional dilaton via flux potential. The point is that the fixed values for dilaton and Kähler moduli from flux compactifications are not consistent with those from the interesting D-brane models and predict the wrong gauge couplings at the string scale for the other models [13]. This is a blessing in disguise from a cosmological point of view [14].

Type IIB Flux Compactifications – We consider the Type IIB string theory compactified on a $\mathbf{T^6}$ orientifold where $\mathbf{T^6}$ is a six-torus factorized as $\mathbf{T^6} = \mathbf{T^2} \times \mathbf{T^2} \times \mathbf{T^2}$ whose complex coordinates are z_i, $i = 1,\ 2,\ 3$ for the i^{th} two-torus, respectively. The orientifold projection is implemented by gauging the symmetry ΩR, where Ω is world-sheet parity, and R is given by

$$R : (z_1, z_2, z_3) \to (-z_1, -z_2, -z_3)\ . \tag{1}$$

Thus, the model contains 64 O3-planes. In order to cancel the negative RR charges from these O3-planes, we introduce the magnetized D(3+2n)-branes which are filling up the four-dimensional Minkowski space-time and wrapping 2n-cycles on the compact manifold. Concretely, for one stack of N_a D-branes wrapped m_a^i times on the i^{th} two-torus $\mathbf{T_i^2}$, we turn on n_a^i units of magnetic fluxes F_a^i for the center of mass $U(1)_a$ gauge factor on $\mathbf{T_i^2}$, such

TABLE I: General spectrum for magnetized D-branes on the Type IIB $\mathbf{T^6}$ orientifold.

Sector	Representation
aa	$U(N_a)$ vector multiplet 3 adjoint multiplets
$ab+ba$	I_{ab} $(N_a, \overline{N}_b)$ multiplets
$ab'+b'a$	$I_{ab'}$ (N_a, N_b) multiplets
$aa'+a'a$	$\frac{1}{2}(I_{aa'} - I_{aO3})$ symmetric multiplets $\frac{1}{2}(I_{aa'} + I_{aO3})$ anti-symmetric multiplets

that

$$m_a^i \, \frac{1}{2\pi} \int_{T_i^2} F_a^i \; = \; n_a^i \; , \tag{2}$$

where m_a^i can be half integer for tilted two-torus. Then, the D9-, D7-, D5- and D3-branes contain 0, 1, 2 and 3 vanishing m_a^is, respectively. Introducing for the i^{th} two-torus the even homology classes $[\mathbf{0}_i]$ and $[\mathbf{T}_i^2]$ for the point and two-torus, respectively, the vectors of the RR charges of the a^{th} stack of D-branes and its image are

$$[\Pi_a] \; = \; \prod_{i=1}^{3} (n_a^i [\mathbf{0}_i] + m_a^i [\mathbf{T}_i^2]),$$
$$[\Pi_a'] \; = \; \prod_{i=1}^{3} (n_a^i [\mathbf{0}_i] - m_a^i [\mathbf{T}_i^2]) \; , \tag{3}$$

respectively. The "intersection numbers" in Type IIA language, which determine the chiral massless spectrum, are

$$I_{ab} \;\; = \;\; [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^{3} (n_a^i m_b^i - n_b^i m_a^i) \; . \tag{4}$$

Moreover, for a stack of N D(2n+3)-branes whose homology classes on $\mathbf{T^6}$ is (not) invariant under ΩR, we obtain a $USp(2N)$ $(U(N))$ gauge symmetry with three anti-symmetric (adjoint) chiral superfields due to the orbifold projection. The physical spectrum is presented in Table I.

The flux models on Type IIB orientifolds with four-dimensional $N = 1$ supersymmetry are primarily constrained by the RR tadpole cancellation conditions that will be given later, the four-dimensional $N = 1$ supersymmetric D-brane configurations, and the K-theory anomaly free conditions. For the D-branes with world-volume magnetic field $F_a^i = n_a^i / (m_a^i \chi_i)$ where χ_i is the area of $\mathbf{T_i^2}$ in string units, the condition for the four-dimensional $N = 1$ supersymmetric D-brane configurations is

$$\sum_i \left(\tan^{-1} (F_a^i)^{-1} + \theta(n_a^i) \pi \right) = 0 \mod 2\pi \; , \tag{5}$$

where $\theta(n_a^i) = 1$ for $n_a^i < 0$ and $\theta(n_a^i) = 0$ for $n_a^i \geq 0$. The K-theory anomaly free conditions are

$$\sum_a N_a m_a^1 m_a^2 m_a^3 = \sum_a N_a m_a^1 n_a^2 n_a^3 = \sum_a N_a n_a^1 m_a^2 n_a^3$$
$$= \sum_a N_a n_a^1 n_a^2 m_a^3 = 0 \mod 2 \; . \tag{6}$$

And the holomorphic gauge kinetic function for a generic stack of D(2n+3)-branes is given by [13, 15]

$$f_a \;\; = \;\; \frac{1}{\kappa_a} \left(n_a^1 \, n_a^2 \, n_a^3 \, s - n_a^1 \, m_a^2 \, m_a^3 \, t_1 \right.$$
$$\left. - n_a^2 \, m_a^1 \, m_a^3 \, t_2 - n_a^3 \, m_a^1 \, m_a^2 \, t_3 \right) \; , \tag{7}$$

where κ_a is equal to 1 and 2 for $U(n)$ and $USp(2n)$, respectively.

We turn on the NSNS flux h_0, RR flux e_i, non-geometric fluxes b_{ii} and $\bar{b}_{ii}$, and the S-dual fluxes f_i, g_{ij} and g_{ii} [12]. To avoid the subtleties, these fluxes should be even integers due to the Dirac quantization. For simplicity, we assume

$$e_i = e \; , \quad b_{ii} = \beta \; , \quad \bar{b}_{ii} = \bar{\beta} \; ,$$
$$f_i = f \; , \; g_{ij} = -g_{ii} = g \; , \tag{8}$$

where $i \neq j$. Then the constraint on fluxes from Bianchi indetities is

$$f \bar{\beta} \;\; = \;\; g \beta \; . \tag{9}$$

The RR tadpole cancellation conditions are

$$\sum_a N_a n_a^1 n_a^2 n_a^3 = 16 \; ,$$
$$\sum_a N_a n_a^i m_a^j m_a^k = -\frac{1}{2} e \bar{\beta} \; ,$$
$$N_{\mathrm{NS7}_i} = 0 \; , \quad N_{\mathrm{I7}_i} = 0 \; , \tag{10}$$

where $i \neq j \neq k \neq i$, and the N_{NS7_i} and N_{I7_i} denote the NS7 brane charge and the other 7-brane charge, respectively [12]. Thus, if $e\bar{\beta} < 0$, the RR tadpole cancellation conditions are relaxed elegantly because $-e\bar{\beta}/2$ only needs to be even integer. Moreover, we have 7 moduli fields in the supergravity theory basis, the dilaton s, three Kähler moduli t_i, and three complex structure moduli u_i. With the above fluxes, we can assume

$$t \equiv t_1 + t_2 + t_3 \; , \quad u_1 = u_2 = u_3 \equiv u \; . \tag{11}$$

Then the superpotential becomes

$$\mathcal{W} = 3ieu + ih_0 s - t \left(\beta u - i \bar{\beta} u^2 \right) - st \left(f - igu \right) . \tag{12}$$

The Kähler potential for these moduli is

$$\mathcal{K} = -\ln(s + \bar{s}) - \sum_{i=1}^{3} \ln(t_i + \bar{t}_i) - \sum_{i=1}^{3} \ln(u_i + \bar{u}_i) \; . \tag{13}$$

In addition, the supergravity scalar potential is

$$V = e^{\mathcal{K}} \left(\mathcal{K}^{i\bar{j}} D_i \mathcal{W} D_{\bar{j}} \mathcal{W} - 3|\mathcal{W}|^2 \right) , \tag{14}$$

where $\mathcal{K}^{i\bar{j}}$ is the inverse metric of $\mathcal{K}_{i\bar{j}} \equiv \partial_i \partial_{\bar{j}} \mathcal{K}$, $D_i\mathcal{W} = \partial_i \mathcal{W} + (\partial_i \mathcal{K})\mathcal{W}$, and $\partial_i = \partial_{\phi_i}$ where ϕ_i can be s, t_i, and u_i. Thus, for the supersymmetric Minkowski vacua, we have

$$\mathcal{W} = \partial_s \mathcal{W} = \partial_t \mathcal{W} = \partial_u \mathcal{W} = 0 . \tag{15}$$

From $\partial_s \mathcal{W} = \partial_t \mathcal{W} = 0$, we obtain

$$t = \frac{ih_0}{f - igu} , \quad s = -\frac{\beta}{f} u , \tag{16}$$

then the superpotential turns out

$$\mathcal{W} = \left(3e - \frac{h_0 \beta}{f} \right) iu . \tag{17}$$

Therefore, to satisfy $\mathcal{W} = \partial_u \mathcal{W} = 0$, we obtain

$$3ef = \beta h_0 . \tag{18}$$

Because $\mathrm{Re}s > 0$, $\mathrm{Re}t_i > 0$ and $\mathrm{Re}u_i > 0$, we require

$$\frac{h_0}{g} < 0 , \quad \frac{\beta}{f} < 0 . \tag{19}$$

Model – Choosing $e\bar{\beta} = -12$, we present the D-brane configurations and intersection numbers in Table II, and the resulting spectrum in Table III. The anomalies from three global $U(1)$s of $U(4)_C$, $U(2)_L$ and $U(2)_R$ are cancelled by the Green-Schwarz mechanism, and the gauge fields of these $U(1)$s obtain masses via the linear $B \wedge F$ couplings. So, the effective gauge symmetry is $SU(4)_C \times SU(2)_L \times SU(2)_R$. In order to break the gauge symmetry, on the first two-torus, we split the a stack of D-branes into a_1 and a_2 stacks with 3 and 1 D-branes, respectively, and split the c stack of D-branes into c_1 and c_2 stacks with 1 D-brane for each one. Then, the gauge symmetry is further broken down to $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$. We can break the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry down to the $U(1)_Y$ gauge symmetry by giving vacuum expectation values (VEVs) to the vector-like particles with quantum numbers $(\mathbf{1}, \mathbf{1}, \mathbf{1/2}, \mathbf{-1})$ and $(\mathbf{1}, \mathbf{1}, \mathbf{-1/2}, \mathbf{1})$ under $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ from $a_2 c_1'$ D-brane intersections. Similar to the discussions in Ref. [9], we can explain the SM fermion masses and mixings via the Higgs fields H_u^i, H_u', H_d^i and H_d' because all the SM fermions and Higgs fields arise from the intersections on the first torus. To decouple the chiral exotic particles, we assume that the T_R^i and S_R^i obtain VEVs at about the string scale, and their VEVs satisfy the D-flatness $U(1)_R$. The chiral exotic particles can obtain masses via the following superpotential

$$W \supset \frac{1}{M_{\mathrm{St}}} S_R^i S_R^j T_R^k T_R^l + T_R^i X^j X^k , \tag{20}$$

where M_{St} is the string scale, and we neglect the $\mathcal{O}(1)$ coefficients in this paper. In addition, the vector-like particles S_L^i and $\overline{S}_L^i$ in the anti-symmetric representation of $SU(2)_L$ can obtain the VEVs close to the string scale while keeping the D-flatness $U(1)_L$. Thus, we can decouple all the Higgs bidoublets close to the string scale except one pair of the linear combinations of the Higgs doublets for the electroweak symmetry breaking at the low energy by fine-tuning the following superpotential

$$\begin{aligned} W \supset\ & \Phi_i(\overline{S}_L^j \Phi' + S_R^j \overline{\Phi}') + \overline{\Phi}_i (T_R^j \Phi' + S_L^j \overline{\Phi}') \\ & + \frac{1}{M_{\mathrm{St}}} \left(\overline{S}_L^i S_R^j \Phi_k \Phi_l + S_L^i T_R^j \overline{\Phi}_k \overline{\Phi}_l \right. \\ & \left. + \overline{S}_L^i T_R^j \Phi' \Phi' + S_L^i S_R^j \overline{\Phi}' \overline{\Phi}' \right) . \end{aligned} \tag{21}$$

In short, below the string scale, we have the supersymmetric SMs which may have zero, one or a few SM singlets from S_L^i, $\overline{S}_L^i$, and/or S_R^i. And then the low bound on the lightest CP-even Higgs boson mass in the minimal supersymmetric SM can be relaxed if we have the SM singlet(s) at low energy [16].

TABLE II: D-brane configurations and intersection numbers.

	$U(4)_C \times U(2)_L \times U(2)_R \times USp(10)$									
	N	(n^i, m^i)	n_S	n_A	b	b'	c	c'	$O3$
a	4	$(1,0) \times (1,-1/2) \times (1,1)$	0	0	3	0(3)	-3	0(3)	0(1)
b	2	$(1,-3) \times (1,1/2) \times (1,0)$	0	0(6)	-	-	0(6)	0(1)	0(3)
c	2	$(1,3) \times (1,1/2) \times (0,-1)$	-6	6	-	-	-	-	3
$O3$	5	$(1,0) \times (2,0) \times (1,0)$	-	-	$6\chi_1 = \chi_2 = 2\chi_3 = 2$				

TABLE III: The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times USp(10)$.

	Quantum Number	Q_4	Q_{2L}	Q_{2R}	Field
ab	$3 \times (4, \bar{2}, 1, 1)$	1	-1	0	$F_L(Q_L, L_L)$
ac	$3 \times (\bar{4}, 1, 2, 1)$	-1	0	1	$F_R(Q_R, L_R)$
c_S	$6 \times (1, 1, \bar{3}, 1, 1)$	0	0	-2	T_R^i
c_A	$6 \times (1, 1, 1, 1, 1)$	0	0	2	S_R^i
$cO3$	$3 \times (1, 1, 2, 10)$	0	0	1	X^i
ac'	$3 \times (4, 1, 2, 1)$	1	0	1	
	$3 \times (\bar{4}, 1, \bar{2}, 1)$	-1	0	-1	
bc	$6 \times (1, 2, \bar{2}, 1)$	0	1	-1	$\Phi_i(H_u^i, H_d^i)$
	$6 \times (1, \bar{2}, 2, 1)$	0	-1	1	$\overline{\Phi}_i$
bc'	$1 \times (1, 2, 2, 1)$	0	1	1	$\Phi'(H_u', H_d')$
	$1 \times (1, \bar{2}, \bar{2}, 1)$	0	-1	-1	$\overline{\Phi}'$
bb'	$6 \times (1, 1, 1, 1)$	0	2	0	S_L^i
	$6 \times (1, \bar{1}, 1, 1)$	0	-2	0	$\overline{S}_L^i$

Next, we consider the gauge coupling unification and moduli stabilization. The real parts of the dilaton and

Kähler moduli in our model are [13]

$$
\mathrm{Re}s = \frac{\sqrt{6}e^{-\phi_4}}{4\pi} ~,~ \mathrm{Re}t_1 = \frac{\sqrt{6}e^{-\phi_4}}{2\pi} ~,
$$
$$
\mathrm{Re}t_2 = \frac{\sqrt{6}e^{-\phi_4}}{12\pi} ~,~ \mathrm{Re}t_3 = \frac{\sqrt{6}e^{-\phi_4}}{6\pi} ~, \qquad (22)
$$

where ϕ_4 is the four-dimensional dilaton. From Eq. (7), we obtain that the SM gauge couplings are unified at the string scale as follows

$$
g^{-2}_{SU(3)_C} = g^{-2}_{SU(2)_L} = \frac{3}{5} g^{-2}_{U(1)_Y} = \frac{\sqrt{6}e^{-\phi_4}}{2\pi} ~. \qquad (23)
$$

Using the unified gauge coupling $g^2 \simeq 0.513$ in supersymmetric SMs, we get

$$
\phi_4 \simeq -1.61 ~. \qquad (24)
$$

For moduli stabilization, we first obtain t from Eqs. (16) and (22)

$$
\mathrm{Re}t = \frac{3\sqrt{6}e^{-\phi_4}}{4\pi} ~,~ \mathrm{Im}t = \pm \sqrt{\frac{3\beta h_0}{fg} - \frac{27 e^{-2\phi_4}}{8\pi^2}} ~. \qquad (25)
$$

Thus, we have

$$
\mathrm{Im}s = -\frac{1}{3}\mathrm{Im}t + \frac{\beta}{g} ~,
$$
$$
\mathrm{Re}u = -\frac{\sqrt{6} f e^{-\phi_4}}{4\pi\beta} ~,~~ \mathrm{Im}u = \frac{f}{3\beta}\mathrm{Im}t - \frac{f}{g} ~. \qquad (26)
$$

Let us present a set of possible solutions to the fluxes

$$
h_0 = -18\eta ~,~~ e = 6\eta ~,~~ \beta = 2\eta' ~,
$$
$$
\bar{\beta} = -2\eta ~,~~ f = -2\eta' ~,~~ g = 2\eta ~, \qquad (27)
$$

where $\eta = \pm 1$ and $\eta' = \pm 1$. Choosing $\phi_4 = -1.61$, $\eta = \eta' = 1$, we obtain the numerical values for the moduli fields

$$
\mathrm{Re}s = \mathrm{Re}u = 0.975 ~,~~~ \mathrm{Re}t_1 = 1.95 ~,
$$
$$
\mathrm{Re}t_2 = 0.325 ~,~~~ \mathrm{Re}t_3 = 0.650 ~,
$$
$$
\sum_{i=1}^{3} \mathrm{Im}t_i = \pm 4.30 ~,~~~ \mathrm{Im}s = \mathrm{Im}u = \mp 1.43 + 1 ~. (28)
$$

Conclusions – We showed that the RR tadpole cancellation conditions can be relaxed elegantly in the supersymmetric Minkowski vacua on the Type IIB toroidal orientifold with general flux compactifications. And we presented a realistic Pati-Salam like model in details. In this model, we can break the gauge symmetry down to the SM gauge symmetry, realize the gauge coupling unification, and decouple all the extra chiral exotic particles around the string scale. We can also generate the observed SM fermion masses and mixings. Futhermore, the unified gauge coupling, the dilaton, the complex structure moduli, the real parts of the Kähler moduli and the sum of the imaginary parts of the Kähler moduli can be determined as functions of the four-dimensional dilaton and fluxes, and can also be estimated.

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